Particle production by time–dependent meson fields in relativistic heavy–ion collisions

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Abstract

According to the Walecka mean–field theory of nuclear interaction the collective mutual deceleration of the colliding nuclei gives rise to the bremsstrahlung of real and virtual ω–mesons. It is shown that decays of these mesons may give a noticeable contribution to the observed yields of the baryon–antibaryon pairs, dileptons and pions. Excitation functions and rapidity distributions of particles produced by this mechanism are calculated under some simplifying assumptions about the space–time variation of meson fields in nuclear collisions. The calculated multiplicities of coherently produced particles grow fast with the bombarding energy, reaching a saturation above the RHIC bombarding energy. In the case of central Au+Au collisions the bremsstrahlung mechanism becomes comparable with particle production in incoherent hadron–hadron collisions above the AGS energies. The rapidity spectra of antibaryons and pions exhibit a characteristic two–hump structure which is a consequence of incomplete projectile–target stopping at the initial stage of the reaction. The predicted distribution of e^+e^- pairs has a strong peak at invariant masses $M_{e^+e^-} < 0.5$ GeV.

1 Introduction

As follows from the relativistic mean–field model [1] strong time–dependent meson fields are generated in the course of relativistic heavy–ion collisions. Within the framework of this model several new collective phenomena were predicted: the filamentation instability of interpenetrating nuclei [2] and the spontaneous creation of the baryon–antibaryon ($B\overline{B}$) pairs in a superdense baryon–rich matter [3]. Using the approach developed in papers on the pion [4] and photon [5] bremsstrahlung we suggested recently [6] a new mechanism of the $B\overline{B}$ pair production by the collective bremsstrahlung of meson fields in relativistic heavy–ion collisions. These pairs may be produced at sufficiently high bombarding energies when characteristic Fourier frequencies of meson fields exceed the energy gap between the positive and negative energy levels of baryons.

In the lowest order approximation the production of the $B\overline{B}$ pair may be considered as a two–step process $A_pA_t \to \omega_* \to B\overline{B}$. Here $A_p(A_t)$ denotes the projectile (target) nucleus and ω_* indicates the off–mass–shell vector meson [1]. The first step in the above reaction is the virtual bremsstrahlung producing virtual mesons with masses $M > 2m_B$ (m_B is the baryon mass). The second step is the conversion of a vector meson into $B\overline{B}$ pairs. Such a process is suppressed for a "real" ω meson which has the mass $m_\omega \simeq 0.783$ GeV and the relatively small width $\Gamma_\omega \simeq 8.4$ MeV. It is clear that analogous bremsstrahlung mechanism may produce also pions (by decays of quasireal mesons $\omega \to \pi^+\pi^0\pi^-$) and low mass dileptons ($M_{l^+l^-} \lesssim m_\omega$).

In this work we study the bremsstrahlung of vector meson fields originated from the collective deceleration of the projectile and target nuclei at the initial stage of a heavy–ion collision. The various channels of the bremsstrahlung conversion, including the production of the $N\overline{N}$ pairs, pions and dileptons are considered with emphasize to their observable signals.

2 Particle production by bremsstrahlung of nuclear meson fields

By the analogy to the Walecka model we introduce the vector meson field $\omega^\mu(x)$ coupled to the 4–current $J^\mu(x)$ of nucleons participating in a heavy–ion collision at a given impact parameter. The equation of motion defining the space–time behavior of $\omega^\mu(x)$ may be written as ($c = \hbar = 1$)

$$(\partial^\nu\partial_\nu + m_\omega^2)\,\omega^\mu(x) = g_V\,J^\mu(x)\,, \tag{1}$$

where g_V is the ωN coupling constant. In the mean–field approximation the quantum fluctuation of J^μ are disregarded and the vector meson field is purely classical. From Eq. (1) one

[1] As discussed in Ref. [6], at relativistic bombarding energies bremsstrahlung of the scalar meson field is small as compared to the vector meson field. Due to this reason we disregard here the contribution of the scalar meson bremsstrahlung.

can see that excitation of propagating waves in a vacuum (bremsstrahlung) is possible if the Fourier transformed baryonic current

$$J^{\mu}(p) = \int \mathrm{d}^4 x J^{\mu}(x) e^{ipx} \tag{2}$$

is nonzero in the time–like region $p^2 = m_{\omega}^2$.

In the following we study the bremsstrahlung process in the lowest order approximation neglecting the back reaction and reabsorption of the emitted vector mesons, i.e. treating J^{μ} as an external current. From Eq. (1) one can calculate the energy flux of the vector field at a large distance from the collision region [4]. This leads to the following formulae for the momentum distribution of real ω–mesons emitted in a heavy–ion collision [2]

$$E_{\omega} \frac{\mathrm{d}^3 N_{\omega}}{\mathrm{d}^3 p} = S(E_{\omega}, \boldsymbol{p}) \,, \tag{3}$$

where $E_{\omega} = \sqrt{m_{\omega}^2 + \boldsymbol{p}^2}$ and

$$S(p) = \frac{g_V^2}{16\pi^3} |J_{\mu}^{*}(p) J^{\mu}(p)| \tag{4}$$

is a source function. In our model the latter is fully determined by the collective motion of the projectile and target nucleons.

To take into account the off–mass–shell effects we characterize virtual ω mesons by their mass M and total width $\Gamma_{\omega *}$. The spectral function describing the deviation from the on–mass–shell may be written as

$$\rho(M) = \frac{2}{\pi} \frac{M \Gamma_{\omega *}}{(M^2 - m_{\omega}^2)^2 + m_{\omega}^2 \Gamma_{\omega *}^2} \,. \tag{5}$$

To calculate the distribution of virtual mesons in their 4–momentum p we use the formulae [7]

$$\frac{\mathrm{d}^4 N_{\omega *}}{\mathrm{d}^4 p} = \rho(M) S(p) \,, \tag{6}$$

where $M \equiv \sqrt{p^2}$. In the limit $\Gamma_{\omega *} \to 0$ one can replace $\rho(M)$ by $2\delta(M^2 - m_{\omega}^2)$. In this case Eq. (6) becomes equivalent to the formulae (3) for the spectrum of the on–mass–shell vector mesons.

Below we consider the most important channels of the virtual ω decay: $i = 3\pi, N\overline{N}$, e^+e^-, $\mu^+\mu^-$. The total width $\Gamma_{\omega *}$ can be decomposed into the sum over the partial decay widths $\Gamma(\omega_* \to i)$:

$$\Gamma_{\omega *} = \sum_i \Gamma(\omega_* \to i)\,. \tag{7}$$

The distribution over the total 4–momentum of particles in a given decay channel may be written as

$$\frac{\mathrm{d}^4 N_{\omega * \to i}}{\mathrm{d}^4 p} = B(\omega_* \to i) \frac{\mathrm{d}^4 N_{\omega *}}{\mathrm{d}^4 p}\,, \tag{8}$$

where $B(\omega_* \to i) \equiv \Gamma(\omega_* \to i)/\Gamma_{\omega *}$ is the branching ratio of the i–th decay channel. The latter is a function of the invariant mass of the decay particles M.

To calculate the 4–vectors $J^\mu(p)$ defining the source function $S(p)$ we assume the simple picture of a high–energy heavy–ion collision suggested in Ref. [6] Below we consider collisions of identical nuclei ($A_p = A_t = A$) at zero impact parameter. In the equal velocity frame the projectile and target nuclei initially move to each other with the velocity $v_0 = (1 - 4m_N^2/s)^{1/2}$, where $\sqrt{s}$ is the c.m. bombarding energy per nucleon. In the "frozen density" approximation [6] the internal compression and transverse motion of nuclear matter are disregarded at the early (interpenetration) stage of the reaction. Within this approximation the colliding nuclei move as a whole along the beam axis with instantaneous velocities $\dot{\xi}_p = -\dot{\xi}_t \equiv \xi(t)$. The projectile velocity $\dot{\xi}(t)$ is a decreasing function of time, chosen in the form [4]

$$\dot{\xi}(t) = v_f + \frac{v_0 - v_f}{1 + \mathrm{e}^{t/\tau}}\,, \tag{9}$$

where τ is the effective deceleration time and v_f is the final velocity of nuclei (at $t \to +\infty$) .

In our approximation the Fourier transforms $J^\mu(p)$ are totally determined by the projectile trajectory $\xi(t)$ [6]:

$$J^0(p) = \frac{p_\parallel}{p_0} J^3(p) = 2A \int\limits_{-\infty}^{\infty} \mathrm{d}t e^{ip_0 t} \cos\left[p_\parallel \xi(t)\right] F\left(\sqrt{\boldsymbol{p}_T^2 + p_\parallel^2 \cdot [1 - \dot{\xi}^2(t)]}\right)\,, \tag{10}$$

where $p_\parallel$ and $\boldsymbol{p}_T$ are, respectively, the longitudinal and transverse components of the three–momentum $\boldsymbol{p}$, $F(q)$ is the density form factor of the initial nuclei

$$F(q) \equiv \frac{1}{A}\int \mathrm{d}^3 r\rho(r)e^{-i\boldsymbol{q}\,\boldsymbol{r}}\,. \tag{11}$$

The time integrals in Eq. (10) were calculated numerically assuming the Woods–Saxon distribution of the nuclear density $\rho(r)$. According to Eqs. (4), (10) the source function $S(p)$ vanishes at $p_\parallel = 0$. As a result, at high bombarding energies single particle distributions have a dip in a central rapidity region (see Figs. 2–3 and Ref. [6]). The two–hump structure of the rapidity spectra is a consequence of the incomplete mutual stopping of nuclei at the initial stage of a heavy–ion collision. On the other hand, the conventional mechanism of particle production in incoherent hadron–hadron collisions results in rapidity distributions of pions and antiprotons with a single central maximum even at high bombarding energies [8, 9].

In this work we use the same choice of the coupling constant g_V and the stopping parameters τ, v_f as in Ref. [6] In particular, it is assumed that τ equals one half of the nuclear passage time

$$\tau = R/\sinh y_0\,, \tag{12}$$

where R is the geometrical radius of initial nuclei. Instead of v_f we introduce the c.m. rapidity loss δy:

$$v_f = \tanh\left(y_0 - \delta y\right). \tag{13}$$

In the case of a central Au+Au collision we assume the energy–independent value [8] $\delta y = 2.4$ for $\sqrt{s} > 10$ GeV and full stopping ($\delta y = y_0$) for lower bombarding energies.

The partial width of the 3π decay channel is calculated assuming that $\Gamma(\omega_* \to 3\pi)$ is proportional to the three–body phase space volume [2]. The normalization constant is determined from the condition that $B(\omega_* \to 3\pi)$ equals the observable value $B(\omega \to 3\pi) = 0.89$ at $M = m_\omega$. Calculation of the $\omega_* \to N\overline{N}$ matrix element in the lowest order approximation in g_V gives the result

$$\Gamma(\omega_* \to N\overline{N}) = \frac{g_V^2}{6\pi}\sqrt{M^2 - 4m_N^2}\left(1 + \frac{2m_N^2}{M^2}\right)\Theta(M - 2m_N)\,, \tag{14}$$

where $\Theta(x) = \frac{1}{2}(1 + \mathrm{sign}\, x)$. After substituting (14) into Eq. (6) and omitting the second term in the denominator of $\rho(M)$ one arrives at the distribution over the pair 4–momentum obtained earlier in Ref. [6]

The dilepton production is studied by calculating the matrix elements of the process $\omega_* \to \gamma_* \to l^+ l^-$ where γ_* is a virtual photon. This calculation leads to the result ($l = \mathrm{e}, \mu$) [10]

$$\frac{\Gamma(\omega_* \to l^+ l^-)}{\Gamma(\omega \to l^+ l^-)} = \left(\frac{m_\omega}{M}\right)^6 \frac{M^2 + 2m_l^2}{m_\omega^2 + 2m_l^2} \sqrt{\frac{M^2 - 4m_l^2}{m_\omega^2 - 4m_l^2}}\, \Theta(M - 2m_l)\,, \tag{15}$$

where m_l is the lepton mass.

3 Results

Below we present the results of numerical calculations obtained within the model described in the preceding section. Some of the model predictions concerning the $B\overline{B}$ production have been already published in Ref. [6]

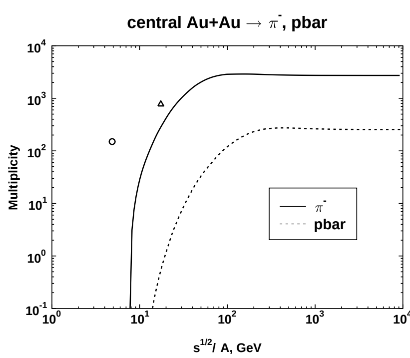

Figure 1: Excitation functions of π^- mesons (solid line) and antiprotons (dashed line) produced by bremsstrahlung in central Au+Au collision. Circle and triangle are experimental data on π^- multiplicity (see text).

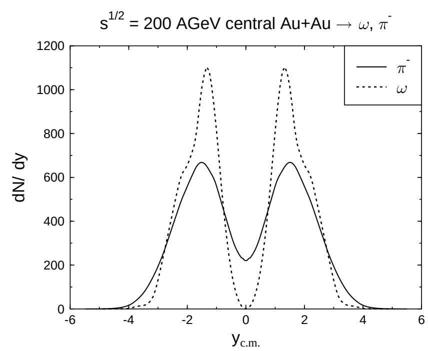

Figure 2: Rapidity distributions of ω (solid line) and π^- (dashed line) mesons produced by bremsstrahlung in central Au+Au collision at RHIC bombarding energy.

Fig. 1 shows the π^- and $\overline{p}$ multiplicities as functions of the bombarding energy in the case of central Au+Au collisions. For comparison, we show experimental data on the π^- multiplicities in the 11.6 AGeV/c Au+Au (circle) [11] and 160 AGeV Pb+Pb (triangle) [12] central collisions. One can see that the multiplicity of pions produced by bremsstrahlung exhibits a rapid growth between the AGS ($\sqrt{s} \simeq 5$ AGeV) and the SPS ($\sqrt{s} \simeq 20$ AGeV) energies and saturates above the RHIC ($\sqrt{s} \simeq 200$ AGeV) energy region. It is interesting that the bremsstrahlung component of pion yield becomes comparable with pion production in incoherent hadron–hadron collisions [9] already at the SPS bombarding energies. Note, however, that actual pion and antiproton yields, especially for heavy combinations of nuclei, may be reduced due to the absorption and annihilation neglected in the present model.

The results on the π^- rapidity spectra are represented in Figs. 2–3. The spectra are calculated in the limit of the on–mass–shell ω mesons, i.e. assuming $\Gamma_{\omega*} = 0$. Here we use the kinematic formulae connecting the pion spectrum and the "primordial" distribution of vector mesons, Eq. (3), suggested in Ref. [2] Similarly to the case of antiprotons [6], the pion rapidity spectrum has a pronounced dip at $y_{\mathrm{c.m.}} \simeq 0$. The two–hump structure of the π and $\overline{p}$ spectra may serve as a signature of the bremsstrahlung mechanism. According to Fig. 3 this structure can be seen only at high enough bombarding energies.

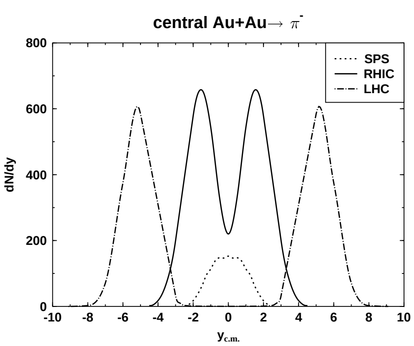

Figure 3: Rapidity spectra of π^- mesons in central Au+Au collisions at different bombarding energies.

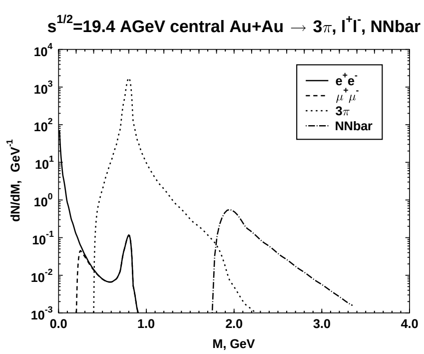

Figure 4: Distributions over invariant masses of particles in different decay channels of virtual ω mesons produced in central Au+Au collision at SPS energy.

Fig. 4 shows the distributions over invariant masses of particles produced in different channels of bremsstrahlung conversion in the case of a central Au+Au collision at the SPS energy. Note that the mass spectrum of e^+e^- pairs created by the bremsstrahlung mechanism has a strong peak at invariant masses below the ω meson mass. On the other hand, attempts to explain the low mass dilepton yield by the conventional incoherent mechanisms (e.g. due to the $\pi\pi \to \rho \to e^+e^-$ processes) strongly underestimate the observable data [13].

4 Conclusions

In this work we have shown that the coherent bremsstrahlung of the vector meson field may be an important source of particle production already at the SPS bombarding energies. The observable signals of this mechanism may be the two–hump structure of pion and antibaryon rapidity spectra as well as the enhanced yield of low mass dileptons. The sharp energy and A–dependence of pion, dilepton and antibaryon excitation functions can be also a signature of the considered mechanism. The latter may be responsible, at least partly, for a rapid increase of the pion multiplicity observed in transition from the AGS to SPS energy [14].

Acknowledgments

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References

[1] B.D. Serot and J.D. Walecka, *Adv. in Nucl. Phys.* **16**, 1 (1985).

[2] Yu.B. Ivanov, *Nucl. Phys.* A **495**, 633 (1989).

[3] I.N. Mishustin, *Yad. Fiz. [Sov. J. Nucl. Phys.]* **52**, 1135 (1990);
I.N. Mishustin, L.M. Satarov, J. Schaffner, H. Stöcker and W. Greiner, *J. Phys. G : Nucl. Part. Phys.* **19**, 1303 (1993).

[4] D. Vasak, H. Stöcker, B. Müller and W. Greiner, *Phys. Lett.* **93B**, 243 (1980);
D. Vasak, B. Müller and W. Greiner, *Phys. Scr.* **22**, 25 (1980).

[5] T. Lippert, U. Becker, N. Grün, W. Scheid and G. Soff, *Phys. Lett.* B **207**, 366 (1988).

[6] I.N. Mishustin, L.M. Satarov, H. Stöcker and W. Greiner, *Phys. Rev.* C **52**, 3323 (1995).

[7] J. Knoll and D.N. Voskresensky, Report GSI-95-63, (Darmstadt, GSI, 1995); *Ann. Phys.* (N.Y.), in print.

[8] Th. Schönfeld, H. Sorge, H. Stöcker and W. Greiner, Mod. Phys. Lett. A **8**, 2631 (1993).

[9] N.S. Amelin, H. Stöcker, W. Greiner, N. Armesto, M.A. Braun and C. Pajares, *Phys. Rev.* C **52**, 362 (1995); N.S. Amelin, private communication.

[10] P. Koch, *Z. Phys.* C **57**, 283 (1993).

[11] M. Gonin, *Nucl. Phys.* A **533**, 799c (1993).

[12] S. Margetis et al., Preprint LBL–36883, (Berkeley, LBL, 1995).

[13] G. Agakishiev et al. (CERES Collab.), *Phys. Rev. Lett.* **75**, 1272 (1995).

[14] M. Gaździcki, D. Röhrich, *Z. Phys.* C **65**, 215 (1995).